

# Signatures of massive sgoldstinos at $e^+e^-$ colliders

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## Abstract

In supersymmetric extensions of the Standard Model with a very light gravitino, the effective theory at the weak scale should contain not only the goldstino  $\tilde{G}$ , but also its supersymmetric partners, the sgoldstinos. In the simplest case, the goldstino is a gauge-singlet and its superpartners are two neutral spin-0 particles,  $S$  and  $P$ . We study possible signals of massive sgoldstinos at  $e^+e^-$  colliders, focusing on those that are most relevant at LEP energies. We show that the LEP constraints on  $e^+e^- \rightarrow \gamma S (\gamma P)$ ,  $ZS (ZP)$  or  $e^+e^- S (e^+e^- P)$ , followed by  $S(P)$  decaying into two gluon jets, can lead to stringent combined bounds on the gravitino and sgoldstino masses.

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# 1. Introduction

A new fundamental scale close to the weak scale,  $G_F^{-1/2} \simeq 300$  GeV, may play a rôle in solving the gauge hierarchy problem of the Standard Model (SM), and allow for unconventional phenomenology at colliders, provided that it can be made compatible with existing data. An old-fashioned example along these lines is technicolor, more fashionable ones identify the new scale with the string scale or some compactification scale. Here we concentrate on a possibility that arises in supersymmetric extensions of the SM, when not only the supersymmetry-breaking mass splittings  $\Delta m^2$ , but also the supersymmetry-breaking scale  $\sqrt{F}$  is close to the weak scale:  $G_F^{-1/2} \sim \Delta m^2 \lesssim \sqrt{F}$ . Since in a flat space-time  $F = \sqrt{3} m_{3/2} M_P$ , where  $m_{3/2}$  is the gravitino mass and  $M_P = (8\pi G_N)^{-1/2} \simeq 2.4 \times 10^{18}$  GeV is the Planck mass, models of this kind are characterized by a very light gravitino, with  $m_{3/2} \lesssim 10^{-3}$  eV.

Many aspects of the superlight gravitino phenomenology at colliders, and in particular gravitino production, either in pairs (tagged by a photon or a jet) or in association with gauginos, have been discussed long ago [1] and also more recently [2, 3]. A very useful tool for these discussions is the supersymmetric equivalence theorem [4], which allows to replace the gravitino with its goldstino components, in the effective theory valid at the present accelerator energies. However, including the goldstino is not the end of the story. If supersymmetry is spontaneously broken but linearly realized in the language of four-dimensional  $N = 1$  local quantum field theory, then the appropriate effective theory must contain, besides the goldstino, also its supersymmetric partners, to be called here *sgoldstinos*.<sup>1</sup> The simplest possible case (as well as the easiest one to reconcile with experimental constraints) corresponds to pure F-breaking of supersymmetry, with the goldstino and the sgoldstinos belonging to a chiral superfield, singlet under the full SM gauge group.

The effective interactions of sgoldstinos with the SM fields and with goldstinos can be characterized by suitable effective couplings. In most cases, at the lowest non-trivial order in a supersymmetric derivative expansion, these couplings are proportional to positive powers of supersymmetry-breaking masses, and to negative powers of the supersymmetry-breaking scale (or, equivalently, of the gravitino mass). Therefore, given the present experimental lower bounds on supersymmetry-breaking masses for particles with SM gauge interactions, sgoldstino production and decay may become phenomenologically relevant at the present collider energies, provided that the sgoldstino masses and the supersymmetry-breaking scale are not too large.

Most of the existing studies of sgoldstino phenomenology [5]–[10] were strongly influenced by the model of ref. [5], where the sgoldstinos were massless at the classical level, and were assumed to acquire very small masses after the inclusion of quantum corrections. As a result, collider signals of sgoldstinos were studied [6]–[9] only in the limit of vanishing

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<sup>1</sup>Notice that, in the presence of an exact R-parity, as assumed throughout this paper, the goldstino is R-odd and the sgoldstinos are R-even.

sgoldstino masses. However, sgoldstino mass terms are allowed by the generic structure of supersymmetric effective Lagrangians, as can be explicitly verified [11]. Moreover, it was recently shown [12] that the situation in which the sgoldstinos and the gravitino are very light, whilst the superpartners of the SM particles are heavy, is generically unstable against quantum corrections, i.e. technically unnatural. Therefore, a more plausible starting point for discussing sgoldstino phenomenology at colliders is to assume that the sgoldstino masses are arbitrary parameters, to be constrained only by experiment. This is the approach that will be followed in the present paper.

The plan of the paper is as follows. In the rest of this section we summarize the assumptions underlying our analysis. These assumptions are much less restrictive than those of the existing literature, but still depart from full generality. Readers who are not familiar with the formalism of supersymmetric effective Lagrangians, and are only interested in the phenomenological aspects of our work, can skip this part and move directly to sect. 2. There we give a systematic discussion of the most important sgoldstino decay modes, focusing on important issues such as the sgoldstino total width and branching ratios. In sect. 3 we discuss the mechanisms for sgoldstino production in  $e^+e^-$  collisions, with emphasis on those that are most important for LEP energies. In sect. 4 we summarize the resulting signals at LEP and we present our conclusions.

We conclude this introduction by recalling the assumptions on the effective theory underlying our analysis. They may be useful for the readers who want to re-derive, starting from the general formalism of supersymmetry, the effective couplings used in the following sections.

We consider a four-dimensional effective theory with global  $N = 1$  supersymmetry and  $SU(3)_C \times SU(2)_L \times U(1)_Y$  gauge symmetry. The building blocks of such a theory are the fields of the Minimal Supersymmetric extension of the Standard Model (MSSM), plus a gauge-singlet chiral superfield  $Z \equiv (z, \tilde{G}, F^Z)$ .

The most general effective Lagrangian with the above field content is determined, up to higher-derivative terms, by a gauge kinetic function  $f$ , a superpotential  $w$  and a Kähler potential  $K$  (see, e.g., ref. [13]). A detailed discussion of the conditions to be imposed on  $f$ ,  $w$  and  $K$  to obtain a fully realistic model will be given elsewhere [14]. Here we give a simplified treatment, mentioning only those features that are directly relevant for the study of sgoldstino phenomenology. First, we make the simplifying assumption that the gauge kinetic function  $f$ , which in principle transforms as a symmetric product of adjoint representations, factorizes into three independent gauge-invariant functions, one for each factor of the gauge group. Then, we assume that the only source of CP violation is the standard Kobayashi-Maskawa phase, so that, apart from the Yukawa couplings, we can restrict ourselves to real parameters and vacuum expectation values (VEVs). As already announced, we assume that supersymmetry is spontaneously broken by  $\langle \partial w / \partial z \rangle \equiv F \neq 0$ , with vanishing VEVs for the auxiliary fields of all the other chiral and vector multiplets. This allows us to identify the fermionic field  $\tilde{G}$  with the goldstino. It is not restrictive to take  $F$  real and positive, so that  $\sqrt{F}$  can be identified with the supersymmetry-breaking scale.

The spin-0 complex field  $z \equiv (S + iP)/\sqrt{2}$  contains then the sgoldstinos, one CP-even ( $S$ ) and the other CP-odd ( $P$ ). We then assume that the gauge symmetry is spontaneously broken by non-vanishing VEVs of the neutral components of the MSSM Higgs doublets,  $H_1$  and  $H_2$ . To guarantee that  $\rho = 1$  at tree level, as in the MSSM, we impose a custodial symmetry on  $K$ , assuming that it depends only on  $(H_1 H_2, \overline{H_1 H_2}, H_1^\dagger H_1 + H_2^\dagger H_2)$ , but not on  $(H_1^\dagger H_1 - H_2^\dagger H_2)$ . To simplify the discussion further, and to avoid the proliferation of parameters, we finally assume that there is no sgoldstino-Higgs mixing, and that squarks, sleptons, gluinos, charginos, neutralinos and Higgs bosons are sufficiently heavy not to play a rôle in sgoldstino production and decay. We can thus take  $S$  and  $P$  as mass eigenstates, with eigenvalues  $m_S^2$  and  $m_P^2$ , respectively. Despite its simplicity, the present context will allow us to generalize considerably the existing studies on sgoldstinos at colliders [6]–[9]. A more general treatment of the interplay between  $SU(2) \times U(1)$  and supersymmetry breaking will be given elsewhere [14].

As explained in [15], we can use the freedom to perform analytic field redefinitions allowed by gauge invariance, and move to a field basis (*normal coordinates*) such that, at the minimum of the potential, all chiral superfields have canonical kinetic terms and, in addition, all the derivatives of the Kähler potential with respect to one chiral superfield and  $n > 1$  antichiral superfields (or vice-versa) have vanishing VEVs. Moreover, this still leaves sufficient freedom to redefine  $Z$  by a suitable constant shift, to ensure that  $\langle z \rangle = 0$ . In the following we shall always assume normal coordinates and  $\langle z \rangle = 0$ : this will lead to simple formulae for the mass spectrum and the interactions, with no further loss of generality.

## 2. Sgoldstino decays

Since we have assumed that squarks, sleptons, gluinos, charginos, neutralinos and Higgs bosons are sufficiently heavy to play no rôle in the decays of the sgoldstinos, we are left with the following possibilities for two-body sgoldstino decays:

$$S(P) \longrightarrow \tilde{G}\tilde{G}, \gamma\gamma, gg, \gamma Z, ZZ, W^+W^-, f\bar{f}, \quad (1)$$

and finally

$$S \longrightarrow PP. \quad (2)$$

Three-body decays such as  $S(P) \rightarrow ggg, W^+W^-\gamma, W^+W^-Z$ , and four-body decays such as  $S(P) \rightarrow gggg, W^+W^-\gamma\gamma, W^+W^-ZZ$  are also possible. We shall neglect all of them here, since they are sub-leading in a perturbative expansion in the gauge coupling constants. We now discuss one by one the different decay channels. Our notation and conventions are the same as in ref. [11]. For simplicity, here and in the following we shall always assume  $m_S, m_P \gg 1$  GeV, to avoid all theoretical subtleties connected with the non-perturbative aspects of the strong interactions: sgoldstinos with masses up to a few GeV would deserve a dedicated phenomenological study.

$$S(P) \longrightarrow \tilde{G}\tilde{G}$$

These decays are controlled by the effective interactions

$$\mathcal{L}_{z\tilde{G}\tilde{G}} = -\frac{1}{2\sqrt{2}}\frac{m_S^2}{F}S\tilde{G}\tilde{G} + \frac{i}{2\sqrt{2}}\frac{m_P^2}{F}P\tilde{G}\tilde{G} + \text{h.c.}, \quad (3)$$

which give [11]

$$\Gamma(S \longrightarrow \tilde{G}\tilde{G}) = \frac{m_S^5}{32\pi F^2}, \quad \Gamma(P \longrightarrow \tilde{G}\tilde{G}) = \frac{m_P^5}{32\pi F^2}. \quad (4)$$

$$S(P) \longrightarrow \gamma\gamma$$

The relevant terms in the effective Lagrangian are

$$\mathcal{L}_{z\gamma\gamma} = -\frac{1}{2\sqrt{2}}\frac{M_{\gamma\gamma}}{F}S F^{\mu\nu}F_{\mu\nu} + \frac{1}{4\sqrt{2}}\frac{M_{\gamma\gamma}}{F}P \epsilon^{\mu\nu\rho\sigma}F_{\mu\nu}F_{\rho\sigma}, \quad (5)$$

where  $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$  is the electromagnetic field strength, and

$$M_{\gamma\gamma} = M_1 \cos^2 \theta_W + M_2 \sin^2 \theta_W. \quad (6)$$

In the above equation,  $M_1$  and  $M_2$  are the diagonal mass terms for the  $U(1)_Y$  and  $SU(2)_L$  gauginos, respectively. From eq. (5) we can easily compute the decay rates, generalizing the results of [11]:

$$\Gamma(S \longrightarrow \gamma\gamma) = \frac{m_S^3 M_{\gamma\gamma}^2}{32\pi F^2}, \quad \Gamma(P \longrightarrow \gamma\gamma) = \frac{m_P^3 M_{\gamma\gamma}^2}{32\pi F^2}, \quad (7)$$

$$S(P) \longrightarrow gg$$

The discussion of these decay modes is a straightforward generalization of the previous ones to the non-Abelian case. They are controlled by the effective interactions

$$\mathcal{L}_{zgg} = -\frac{1}{2\sqrt{2}}\frac{M_3}{F}S G^{\mu\nu\alpha}G_{\mu\nu}^\alpha + \frac{1}{4\sqrt{2}}\frac{M_3}{F}P \epsilon^{\mu\nu\rho\sigma}G_{\mu\nu}^\alpha G_{\rho\sigma}^\alpha, \quad (8)$$

where  $G_{\mu\nu}^\alpha$  is the  $SU(3)$  field strength and  $M_3$  is the gluino mass. From eq. (8) we obtain

$$\Gamma(S \longrightarrow gg) = \frac{m_S^3 M_3^2}{4\pi F^2}, \quad \Gamma(P \longrightarrow gg) = \frac{m_P^3 M_3^2}{4\pi F^2}. \quad (9)$$

$$S(P) \longrightarrow \gamma Z$$

Since the gaugino block of the neutralino mass matrix and the gauge boson mass matrix cannot be simultaneously diagonalized (apart from the special case  $M_2 = M_1$ ), these decay modes may be phenomenologically relevant and cannot be neglected. They are controlled by the effective interactions [14]

$$\mathcal{L}_{z\gamma Z} = -\frac{1}{\sqrt{2}} \frac{M_{\gamma Z}}{F} S F^{\mu\nu} Z_{\mu\nu} + \frac{1}{2\sqrt{2}} \frac{M_{\gamma Z}}{F} P \epsilon^{\mu\nu\rho\sigma} F_{\mu\nu} Z_{\rho\sigma}, \quad (10)$$

where  $Z_{\mu\nu} = \partial_\mu Z_\nu - \partial_\nu Z_\mu$  is the Abelian field strength for the  $Z$  boson, and

$$M_{\gamma Z} = (M_2 - M_1) \sin \theta_W \cos \theta_W. \quad (11)$$

From eq. (10) we can easily compute the decay rates

$$\Gamma(S \longrightarrow \gamma Z) = \frac{M_{\gamma Z}^2 m_S^3}{16\pi F^2} \left(1 - \frac{m_Z^2}{m_S^2}\right)^3, \quad \Gamma(P \longrightarrow \gamma Z) = \frac{M_{\gamma Z}^2 m_P^3}{16\pi F^2} \left(1 - \frac{m_Z^2}{m_P^2}\right)^3. \quad (12)$$

$$S(P) \longrightarrow ZZ$$

The discussion of these decay modes is complicated by the fact that the corresponding interactions originate not only from the generalized kinetic terms for the electroweak gauge bosons, but also, in the case of  $S$ , from the generalized kinetic terms of the Higgs bosons. These decay modes are controlled by the effective interactions [14]

$$\mathcal{L}_{zZZ} = -\frac{1}{2\sqrt{2}} \frac{M_{ZZ}}{F} S Z^{\mu\nu} Z_{\mu\nu} - \frac{m_Z^2 \mu_a}{\sqrt{2}F} S Z^\mu Z_\mu + \frac{1}{4\sqrt{2}} \frac{M_{ZZ}}{F} P \epsilon^{\mu\nu\rho\sigma} Z_{\mu\nu} Z_{\rho\sigma}, \quad (13)$$

where

$$M_{ZZ} = M_1 \sin^2 \theta_W + M_2 \cos^2 \theta_W, \quad (14)$$

and  $\mu_a$  is a diagonal mass term for the antisymmetric neutralino combination,  $(\tilde{H}_1^0 - \tilde{H}_2^0)/\sqrt{2}$ , analogous (but not identical) to the so-called  $\mu$ -term of the MSSM. From eq. (13) we find

$$\begin{aligned} \Gamma(S \longrightarrow ZZ) &= \frac{1}{32\pi F^2 m_S} \left[ M_{ZZ}^2 \left( m_S^4 - 4m_S^2 m_Z^2 + 6m_Z^4 \right) \right. \\ &\quad - 12M_{ZZ} \mu_a m_Z^2 \left( \frac{m_S^2}{2} - m_Z^2 \right) \\ &\quad \left. + 2\mu_a^2 m_Z^4 \left( \frac{m_S^4}{4m_Z^4} - \frac{m_S^2}{m_Z^2} + 3 \right) \right] \sqrt{1 - \frac{4m_Z^2}{m_S^2}}, \end{aligned} \quad (15)$$

$$\Gamma(P \longrightarrow ZZ) = \frac{M_{ZZ}^2 m_P^3}{32\pi F^2} \left(1 - \frac{4m_Z^2}{m_P^2}\right)^{3/2}. \quad (16)$$

$$S(P) \longrightarrow W^+ W^-$$

The discussion of these decay modes is the obvious generalization of the previous ones. They are controlled by the effective interactions [14]

$$\mathcal{L}_{zWW} = -\frac{1}{\sqrt{2}} \frac{M_2}{F} S W^{\mu\nu} + W_{\mu\nu}^- - \frac{\sqrt{2} m_W^2 \mu_a}{F} S W^{\mu} + W_{\mu}^- + \frac{1}{2\sqrt{2}} \frac{M_2}{F} P \epsilon^{\mu\nu\rho\sigma} W_{\mu\nu}^+ W_{\rho\sigma}^-, \quad (17)$$

where  $W_{\mu\nu}^{\pm} = \partial_{\mu} W_{\nu}^{\pm} - \partial_{\nu} W_{\mu}^{\pm}$  are the Abelian field strengths for the  $W^{\pm}$  bosons, and  $\mu_a$ , already defined above, can be also identified with the diagonal higgsino entry in the chargino mass matrix. From eq. (17) we can easily compute the decay rates

$$\begin{aligned} \Gamma(S \longrightarrow W^+ W^-) &= \frac{1}{16\pi F^2 m_S} \left[ M_2^2 (m_S^4 - 4m_S^2 m_W^2 + 6m_W^4) \right. \\ &\quad - 12M_2 \mu_a m_W^2 \left( \frac{m_S^2}{2} - m_W^2 \right) \\ &\quad \left. + 2\mu_a^2 m_W^4 \left( \frac{m_S^4}{4m_W^4} - \frac{m_S^2}{m_W^2} + 3 \right) \right] \sqrt{1 - \frac{4m_W^2}{m_S^2}}, \end{aligned} \quad (18)$$

$$\Gamma(P \longrightarrow W^+ W^-) = \frac{M_2^2 m_P^3}{16\pi F^2} \left( 1 - \frac{4m_W^2}{m_P^2} \right)^{3/2}. \quad (19)$$

$$S(P) \longrightarrow f \bar{f}$$

As discussed in [16], the Yukawa couplings of sgoldstinos to fermions are expected to be suppressed by a factor  $m_f/\sqrt{F}$ , where  $m_f$  is the fermion mass. This can be justified in terms of the same chiral symmetry that suppresses the off-diagonal (left-right) contributions to the sfermion mass matrices. We then expect these decay modes to be important only for very heavy sgoldstinos decaying into top-antitop pairs, thus we shall always neglect them in the following.

$$S \longrightarrow PP$$

If  $m_S > 2m_P$ , this decay is kinematically allowed and must be taken into account. At the level of the effective theory, the corresponding cubic coupling is a free parameter, unrelated with the spectrum. In the following we shall always neglect this decay mode, consistently with the strategy of focusing the attention on the lighter sgoldstino, the most likely to be discovered first.

Now that we have explicit formulae for the most important two-body decays of the sgoldstinos, we can move to the discussion of their total widths and branching ratios.

Since the expressions for the partial widths of  $S$  and  $P$  are identical in all cases of practical interest, we shall give such a discussion only in the case of  $S$ .

The parameters controlling  $S$  decays are  $m_S$ ,  $\sqrt{F}$ , the gaugino masses ( $M_3, M_2, M_1$ ) and the higgsino mass  $\mu_a$ . In the following, when giving numerical examples, we shall focus our attention on the dependences on  $m_S$  and  $\sqrt{F}$ , by making for the remaining parameters the two representative choices given in table 1. For most values of  $\sqrt{F}$  to be considered in the following, these choices should be comfortably compatible with the present experimental limits on R-odd supersymmetric particles, coming from LEP and Tevatron searches.

	$M_3$	$M_2$	$M_1$	$\mu_a$
(a)	400	300	200	300
(b)	350	350	350	350

Table 1: Two representative choices for the gaugino and higgsino mass parameters affecting sgoldstino decays. All masses are expressed in GeV.

Since all the two-body decay widths are proportional to  $F^{-2}$ , the dependence on  $F$  drops out of the discussion of the  $S$  branching ratios. The latter are shown in fig. 1, as functions of  $m_S$ , for the two parameter choices of table 1. We can see that the dominant decay mode is always the one into gluons. Even in the extreme case (b), this decay mode dominates over the one into photons, because of the color factor 8. Decays into goldstinos can become important only when  $m_S$  is of the order of  $M_3$  or larger: in this case, however, we would expect gauginos (produced in pairs or in association with a gravitino) to be detected before sgoldstinos.

The other important quantity is the total  $S$  width,  $\Gamma_S$ , controlled by the ratios between the relevant mass parameters and the supersymmetry-breaking scale. Small values of these ratios correspond to relatively long-lived sgoldstinos, large values of these ratios correspond to broad, strongly coupled sgoldstinos: to keep the particle interpretation and the validity of our approximations, we must require, among the other things, that  $\Gamma_S/m_S \ll 1$ . We show in fig. 2 contours corresponding to constant values of  $\Gamma_S/m_S$  in the  $(m_S, \sqrt{F})$ -plane, for the parameter choices of table 1. As we shall see in sects. 3 and 4, the region of parameter space of present experimental interest is such that sgoldstinos can always be treated as very narrow resonances.

A question that should be asked is whether the sgoldstino widths can be so small that sgoldstinos can travel for an experimentally significant length in a detector before decaying. Since, as we have seen, the dominant decay mode is always the one into gluons, the typical distance traveled by a sgoldstino  $S$  of mass  $m_S$  and energy  $E_S$  can be written



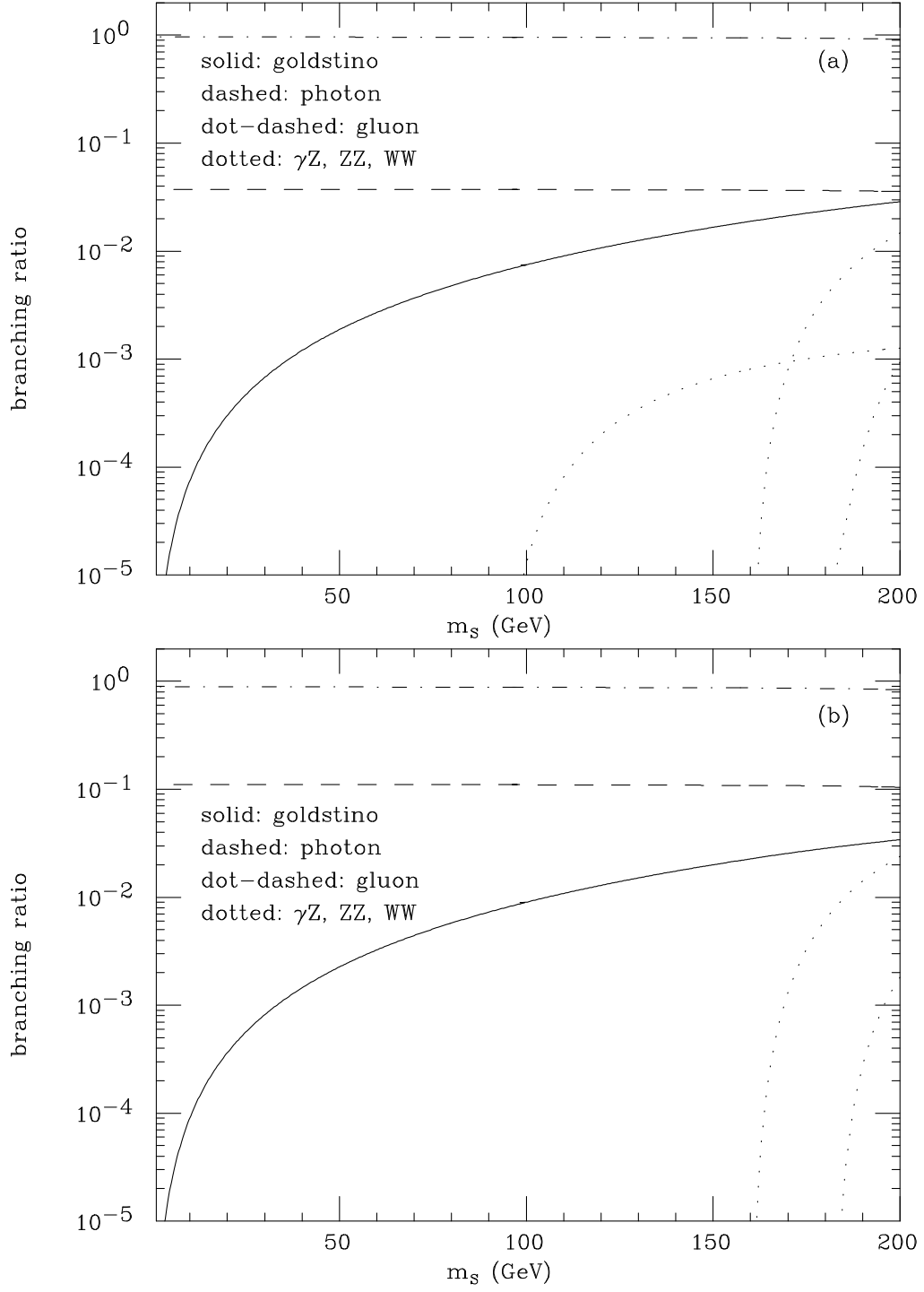


Figure 1: *The  $S$  branching ratios, as functions of  $m_S$ , for the parameter choices of table 1.*

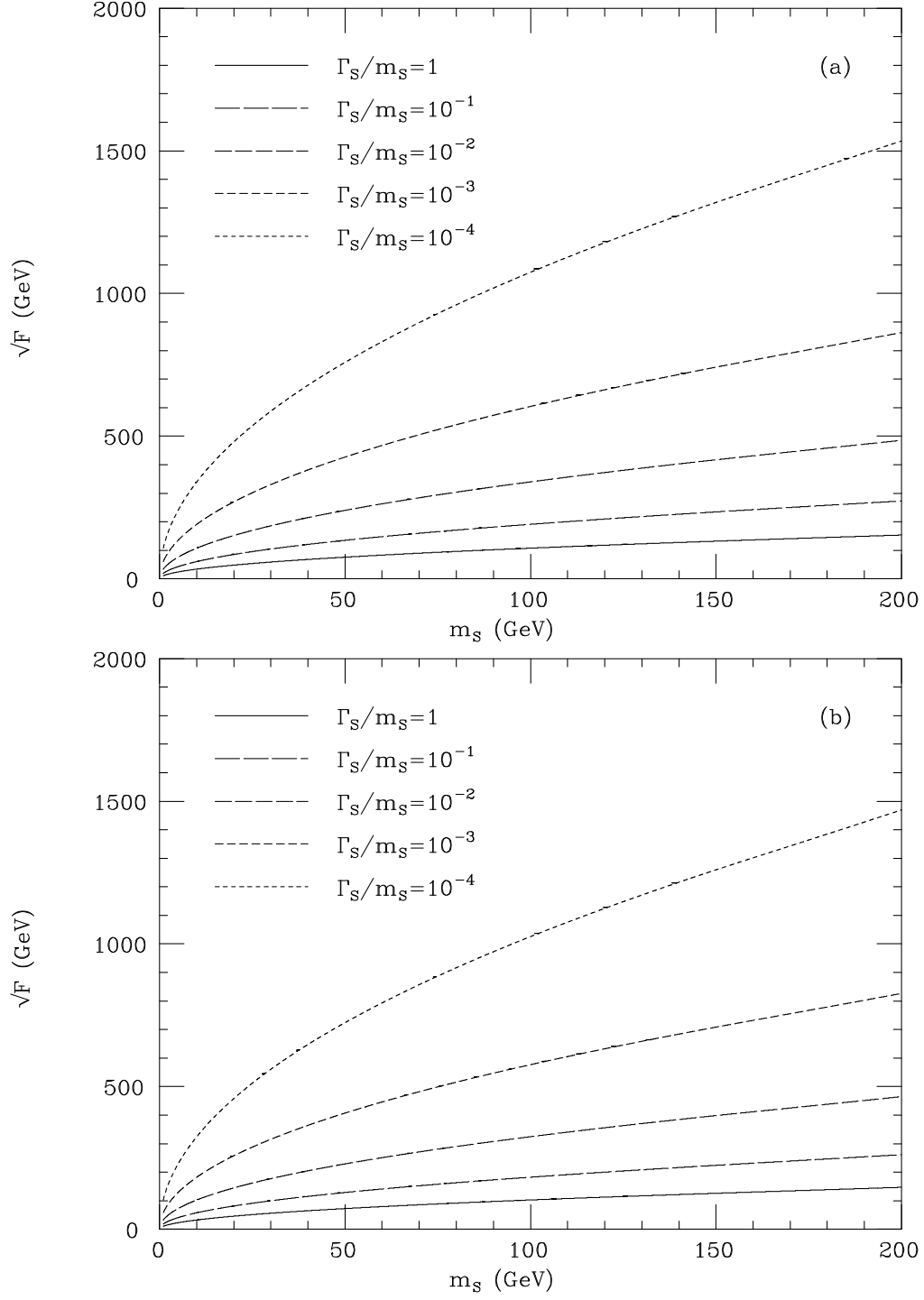


Figure 2: Lines corresponding to fixed values of  $\Gamma_S/m_S$  in the  $(m_S, \sqrt{F})$  plane, for the parameter choices of table 1.

as

$$L \simeq \left( \frac{\sqrt{F}}{1 \text{ TeV}} \right)^4 \left( \frac{1 \text{ GeV}}{m_S} \right)^3 \left( \frac{300 \text{ GeV}}{M_3} \right)^2 \sqrt{\frac{E_S^2}{m_S^2} - 1} \cdot (2.75 \times 10^{-2} \mu m) . \quad (20)$$

Again, we shall check in sects. 3 and 4 that, for  $m_S \gg 1 \text{ GeV}$ , and values of the other parameters not excluded by the present data but leading to non-negligible production cross-sections at LEP, sgoldstinos always decay within a  $\mu m$  from the primary interaction vertex.

Before concluding this section, we should also mention the possibility of three-body decays such as  $S \rightarrow P f \bar{f}$  (or  $P \rightarrow S f \bar{f}$ ), where  $f$  is a light matter fermion, induced by local four-point interactions that are not controlled by the gauge couplings, of the form

$$\mathcal{L}_{f\bar{f}zz} = \frac{1}{2F^2} \left( \tilde{m}_f^2 \bar{f} \bar{\sigma}^\mu f + \tilde{m}_{f^c}^2 \bar{f}^c \bar{\sigma}^\mu f^c \right) (S \partial_\mu P - P \partial_\mu S) , \quad (21)$$

where  $\tilde{m}_f^2$  and  $\tilde{m}_{f^c}^2$  are the diagonal supersymmetry-breaking masses for the left- and right-handed sfermions, respectively. The corresponding widths are easily calculated in the limit of massless fermions:

$$\Gamma(S \rightarrow P f \bar{f}) = \frac{N_f (\tilde{m}_f^4 + \tilde{m}_{f^c}^4)}{6144 \pi^3 m_S^2 F^4} \left[ \frac{m_S^8 - 8 m_S^6 m_P^2 + 8 m_S^2 m_P^6 - m_P^8}{m_S} + 12 m_P^4 m_S^3 \log \frac{m_S^2}{m_P^2} \right] , \quad (22)$$

where  $N_f = 1$  for leptons and  $N_f = 3$  for quarks, and similarly for  $P \rightarrow S f \bar{f}$ . Equation (22) simplifies considerably when the mass of the sgoldstino in the final state can be neglected:

$$\Gamma(S \rightarrow P f \bar{f}) = \frac{N_f (\tilde{m}_f^4 + \tilde{m}_{f^c}^4) m_S^5}{6144 \pi^3 F^4} , \quad (m_P = 0) . \quad (23)$$

Notice, however, that these decays are strongly suppressed not only by the phase space, but also by a higher power of the supersymmetry-breaking scale at the denominator. Moreover, this decay mode is of course relevant only for the heavier sgoldstino, thus, on the same basis as for  $S \rightarrow PP$ , we could safely neglect it in the previous discussion.

### 3. Sgoldstino production

We now review the most important mechanisms for sgoldstino production at  $e^+e^-$  colliders, and especially at LEP. As before, whenever the formulae for  $S$  and  $P$  are identical in form, apart from the obvious substitution  $S \leftrightarrow P$ , we refer to  $S$  only.

Since the sgoldstino couplings to light fermions are suppressed by the corresponding fermion masses (as it is the case for the SM Higgs), resonant production in the  $s$ -channel can be neglected.

At LEPI, we can consider the possibility of  $Z \rightarrow S \gamma$  decays, whose rate can be easily calculated from the effective Lagrangian of eq. (10):

$$\Gamma(Z \rightarrow S \gamma) = \frac{M_{\gamma Z}^2 m_Z^3}{48 \pi F^2} \left( 1 - \frac{m_S^2}{m_Z^2} \right)^3 . \quad (24)$$

To give a measure of the LEPI sensitivity, we show in fig. 3 contours of constant branching ratios,  $BR(Z \rightarrow S\gamma) \equiv \Gamma(Z \rightarrow S\gamma)/\Gamma_Z$ , in the  $(m_S, \sqrt{F})$  plane, for the parameter choice (a) of table 1 [the parameter choice (b) leads to a vanishing effective coupling,  $M_{\gamma Z} = 0$ ].

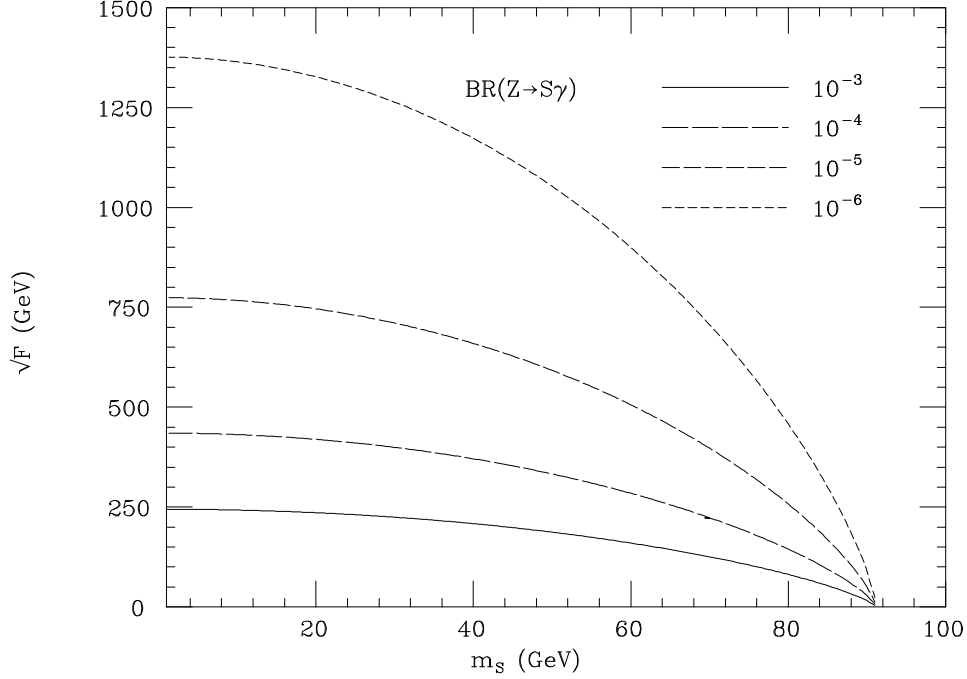


Figure 3: Lines corresponding to fixed values of  $BR(Z \rightarrow S\gamma)$ , in the  $(m_S, \sqrt{F})$  plane, for the parameter choice (a) of table 1.

More generally, we can consider the process  $e^+e^- \rightarrow S\gamma$ . At the classical level, there are only two Feynman diagrams to compute, corresponding to s-channel  $\gamma$  and  $Z$  exchange. Neglecting both the electron mass and the  $Z$  width, the differential cross-section is

$$\frac{d\sigma}{d\cos\theta}(e^+e^- \rightarrow S\gamma) = \frac{|\Sigma|^2 s}{64\pi F^2} \left(1 - \frac{m_S^2}{s}\right)^3 (1 + \cos^2\theta), \quad (25)$$

where

$$|\Sigma|^2 = \frac{e^2 Q_e^2 M_{\gamma\gamma}^2}{2s} + \frac{g_Z^2 (v_e^2 + a_e^2) M_{\gamma Z}^2 s}{2(s - m_Z^2)^2} + \frac{e Q_e g_Z v_e M_{\gamma\gamma} M_{\gamma Z}}{(s - m_Z^2)}, \quad (26)$$

$v_e = T_{3e}/2 - Q_e \sin^2\theta_W$ ,  $a_e = -T_{3e}/2$ ,  $T_{3e} = -1/2$ ,  $Q_e = -1$ ,  $g_Z = e/(\sin\theta_W \cos\theta_W)$ , and  $\theta$  is the scattering angle in the center-of-mass frame. In the limit  $m_S = 0$ ,  $M_1 = M_2$  (i.e.  $M_{\gamma Z} = 0$ ), we recover the results of ref. [8]. This process is particularly relevant at LEP II. To give a measure of the LEP II sensitivity, we draw in fig. 4 contours of constant  $\sigma(e^+e^- \rightarrow S\gamma)$  in the  $(m_S, \sqrt{F})$  plane, for  $\sqrt{s} = 200$  GeV and the two parameter choices of table 1.

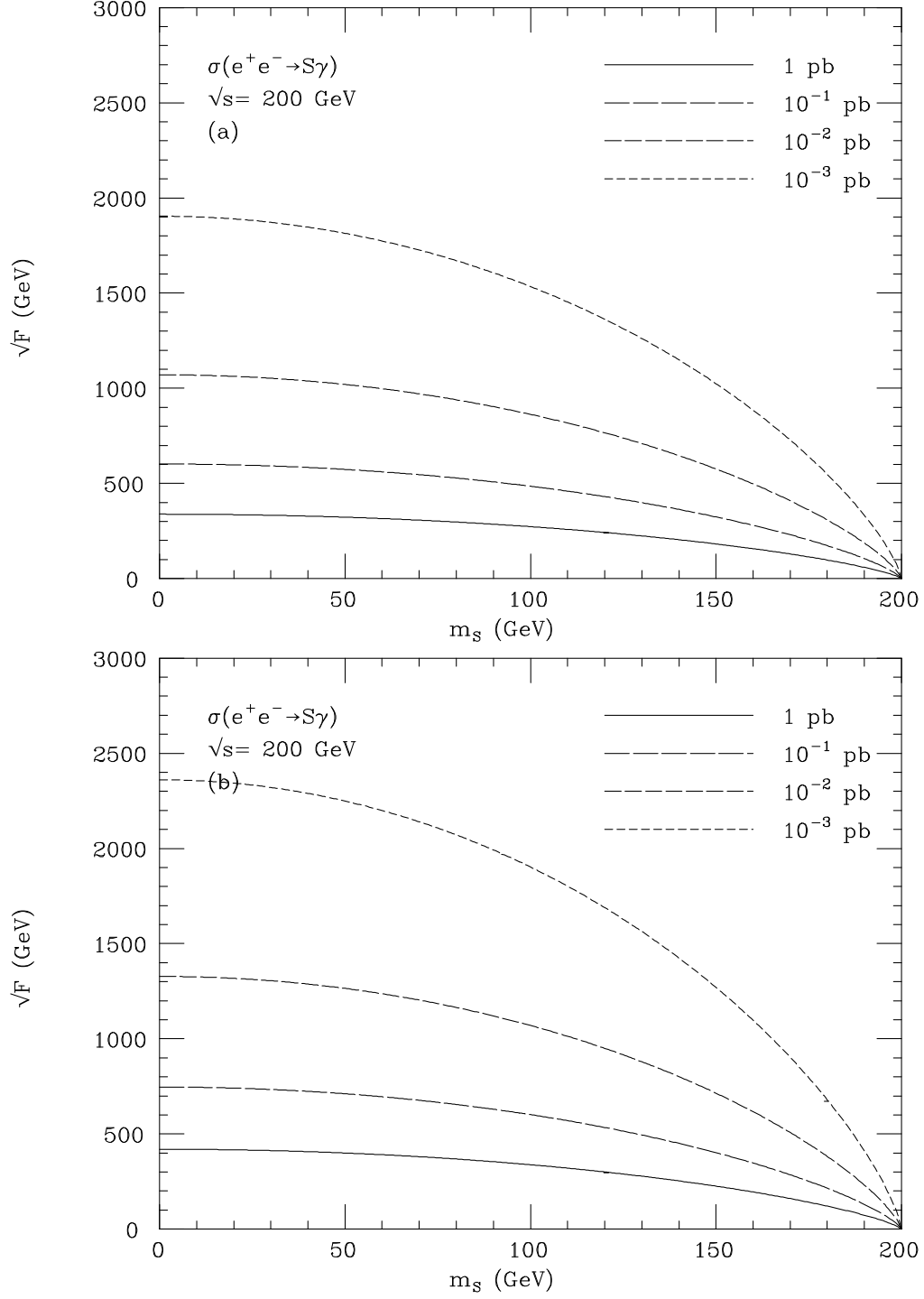


Figure 4: Lines corresponding to fixed values of  $\sigma(e^+e^- \rightarrow S\gamma)$  for  $\sqrt{s} = 200$  GeV, in the  $(m_S, \sqrt{F})$  plane and for the parameter choices of table 1.

Another process to be considered is  $e^+e^- \rightarrow PZ$ , an obvious generalization of  $e^+e^- \rightarrow P\gamma$ . The differential cross-section is given by

$$\frac{d\sigma}{d\cos\theta} (e^+e^- \rightarrow PZ) = \frac{|\Sigma_P|^2}{32\pi s^2 F^2} \sqrt{(s - m_P^2 - m_Z^2)^2 - 4m_P^2 m_Z^2}, \quad (27)$$

where

$$|\Sigma_P|^2 = \left( \frac{e^2 Q_e^2 M_{\gamma Z}^2}{2s} + \frac{g_Z^2 (v_e^2 + a_e^2) M_{ZZ}^2 s}{2(s - m_Z^2)^2} + \frac{e Q_e g_Z v_e M_{\gamma Z} M_{ZZ}}{(s - m_Z^2)} \right) (t^2 + u^2 - 2m_P^2 m_Z^2), \quad (28)$$

and  $t$  and  $u$  are the Mandelstam variables. The cross-section for  $e^+e^- \rightarrow SZ$  has some additional complications, because of the couplings controlled by the higgsino mass parameter  $\mu_a$ :

$$\frac{d\sigma}{d\cos\theta} (e^+e^- \rightarrow SZ) = \frac{|\Sigma_S|^2}{32\pi s^2 F^2} \sqrt{(s - m_S^2 - m_Z^2)^2 - 4m_S^2 m_Z^2}, \quad (29)$$

where

$$\begin{aligned} |\Sigma_S|^2 = & \left[ \frac{e^2 Q_e^2 M_{\gamma Z}^2}{2s} + \frac{g_Z^2 (v_e^2 + a_e^2) M_{ZZ}^2 s}{2(s - m_Z^2)^2} + \frac{e Q_e g_Z v_e M_{\gamma Z} M_{ZZ}}{s - m_Z^2} \right] [t^2 + u^2 + 2m_Z^2 (2s - m_S^2)] \\ & + \frac{g_Z^2 \mu_a^2 m_Z^4 (v_e^2 + a_e^2)}{(s - m_Z^2)^2} \left( 2s - m_S^2 + \frac{tu}{m_Z^2} \right) \\ & + \frac{g_Z \mu_a m_Z^2}{s - m_Z^2} \left[ \frac{g_Z (v_e^2 + a_e^2) M_{ZZ}}{s - m_Z^2} + \frac{e Q_e M_{\gamma Z} v_e}{s} \right] [2s(s + m_Z^2 - m_S^2)]. \end{aligned} \quad (30)$$

We draw in figs. 5 and 6 contours of constant  $\sigma(e^+e^- \rightarrow PZ)$  and  $\sigma(e^+e^- \rightarrow SZ)$ , in the  $(m_P, \sqrt{F})$  and in the  $(m_S, \sqrt{F})$  plane, respectively, for  $\sqrt{s} = 200$  GeV and the two parameter choices of table 1.

Other interesting processes at LEP II are  $e^+e^- \rightarrow e^+e^-S$ , occurring via  $\gamma\gamma$ -,  $\gamma Z$ - and  $ZZ$ -fusion, and  $e^+e^- \rightarrow \nu_e \bar{\nu}_e S$ , occurring via  $WW$ -fusion. We discuss here only the first process, considering only the  $\gamma\gamma$ -fusion diagram, since at LEP energies all the other diagrams give much smaller contributions to the total cross-section, and the interference with  $e^+e^- \rightarrow Z^{(*)}S \rightarrow e^+e^-S$  is negligible. The production of sgoldstinos via  $\gamma\gamma$ -fusion can be described in the Weizsäcker-Williams approximation, i.e. neglecting contributions of off-shell (non-collinear) photons. So doing, the cross-section for the process  $e^+e^- \rightarrow e^+e^-S$  can be expressed in terms of the cross-section for the subprocess  $\gamma\gamma \rightarrow S$ , where each photon is taken on its mass shell, and carries a fraction of the incoming electron momentum which is distributed according to the Weizsäcker-Williams function. In formulae [17]:

$$\sigma(s) = \int_{\tau_S}^1 dx_1 \int_{\tau_S/x_1}^1 dx_2 f^{WW}(x_1) f^{WW}(x_2) d\sigma_{\gamma\gamma}(x_1 x_2 s), \quad (31)$$

where  $s$  is the center-of-mass squared energy,  $\tau_S = m_S^2/s$  and  $x_1, x_2$  are the fractions of the incoming electron and positron momenta carried by the colliding photons. The Weizsäcker-Williams distribution function is given by

$$f^{WW}(x) = \frac{\alpha_{em}}{2\pi} \left[ 2m_e^2 x \left( \frac{1}{m_S^2} - \frac{1-x}{m_e^2 x^2} \right) + \frac{1+(1-x)^2}{x} \log \frac{Q^2(1-x)}{m_e^2 x^2} \right] \quad (32)$$

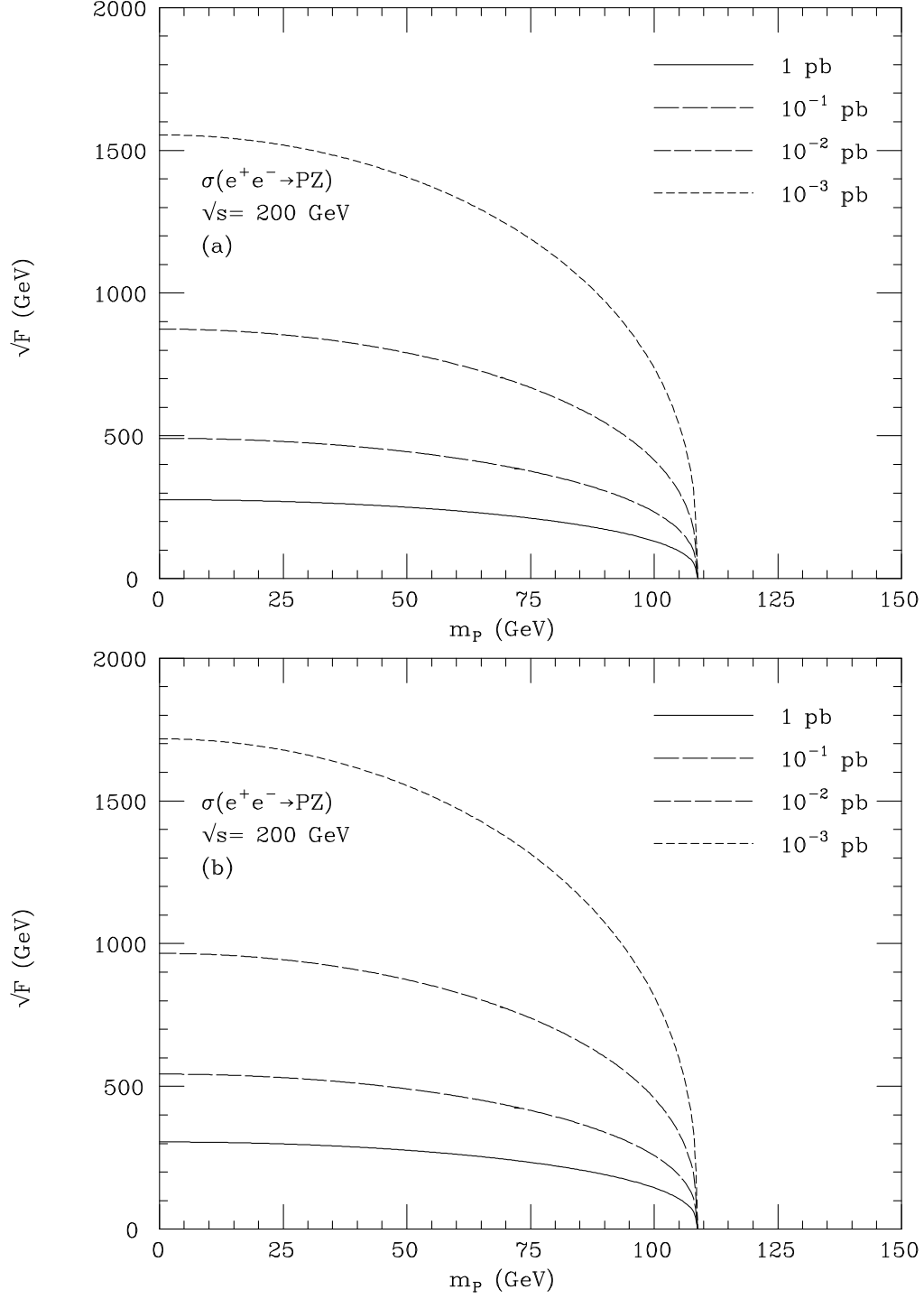


Figure 5: Lines corresponding to fixed values of  $\sigma(e^+e^- \rightarrow PZ)$  for  $\sqrt{s} = 200$  GeV, in the  $(m_P, \sqrt{F})$  plane and for the parameter choices of table 1.

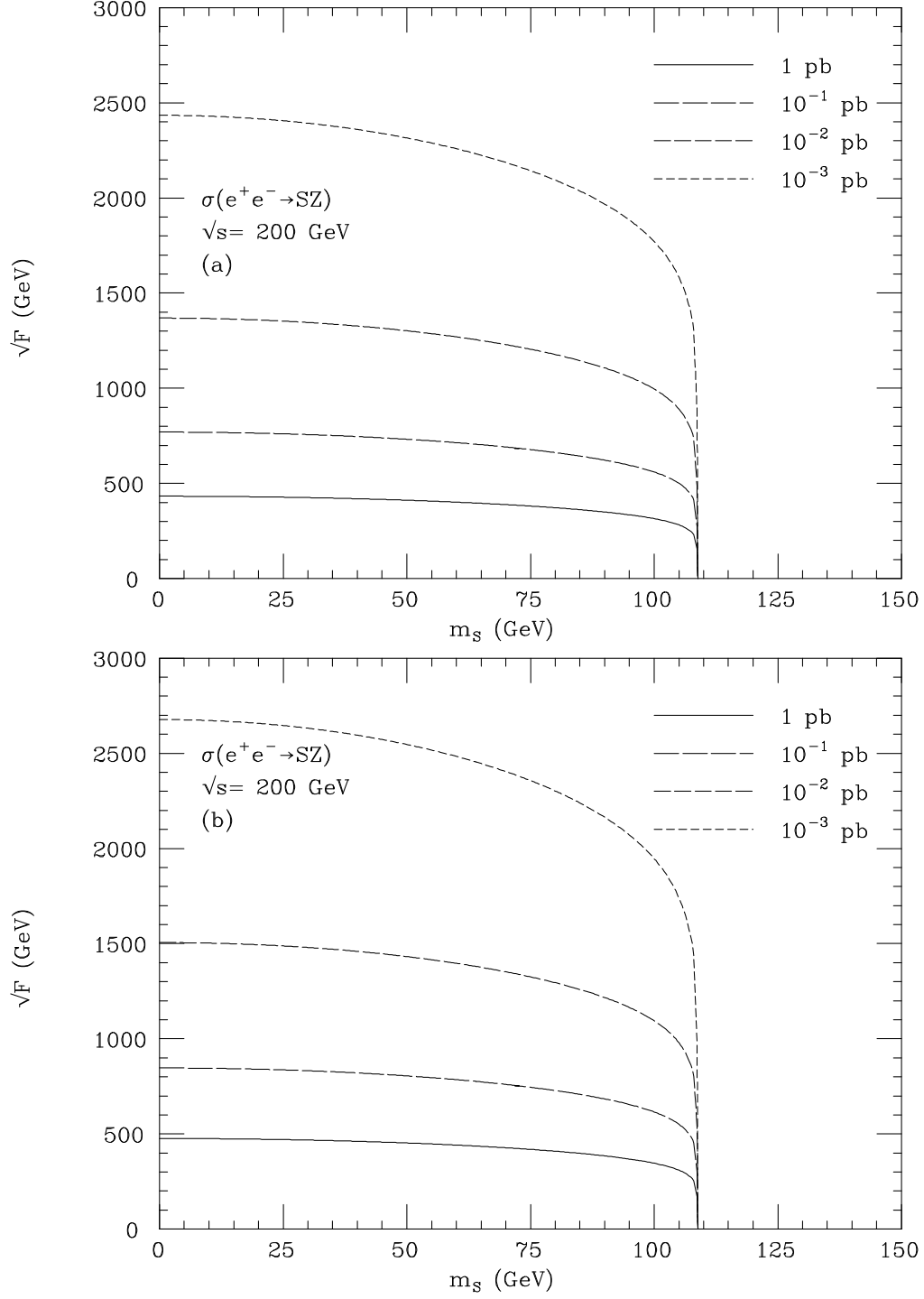


Figure 6: Lines corresponding to fixed values of  $\sigma(e^+e^- \rightarrow S Z)$  for  $\sqrt{s} = 200$  GeV, in the  $(m_S, \sqrt{F})$  plane and for the parameter choices of table 1.



where  $Q$  is an energy scale of the order of  $m_S$ . Corrections to eq. (31) are suppressed with respect to  $\log(m_e/m_S)$  by powers of  $m_e/m_S$  (including the zeroth power). From the effective Lagrangian of eq. (5), we find

$$\sigma_{\gamma\gamma}(\hat{s}) = \frac{M_{\gamma\gamma}^2 m_S^2 \pi}{4F^2} \delta(\hat{s} - m_S^2), \quad (33)$$

from which we deduce, after some trivial manipulations,

$$\sigma(e^+e^- \rightarrow e^+e^-S) = \frac{M_{\gamma\gamma}^2 \tau_S \pi}{4F^2} \int_{\tau_S}^1 \frac{dx}{x} f^{WW}(x) f^{WW}(\tau_S/x). \quad (34)$$

Numerical results are given in fig. 7, which shows contours of constant  $\sigma(e^+e^- \rightarrow e^+e^-S)$  in the  $(m_S, \sqrt{F})$  plane, for  $\sqrt{s} = 200$  GeV and the two parameter choices of table 1.

We conclude this section with a comment similar to the one given at the end of sect. 2. The local effective interaction of eq. (21) can also lead to the pair-production of a CP-even and a CP-odd sgoldstino, with cross-section

$$\frac{d\sigma}{d\cos\theta}(e^+e^- \rightarrow SP) = \frac{(\tilde{m}_e^4 + \tilde{m}_{ec}^4)}{512\pi s^2 F^4} [(s - m_S^2 - m_P^2)^2 - 4m_S^2 m_P^2]^{3/2} \sin^2\theta, \quad (35)$$

where  $\theta$  is the scattering angle in the center-of-mass frame. For plausible values of the parameters, we expect this cross-section to be suppressed by the large numerical factor and the higher power of the supersymmetry-breaking scale at the denominator. Otherwise, the corresponding signal could be seen as an anomaly in the four-jet sample.

#### 4. Discussion of sgoldstino signals and conclusions

The results of sects. 2 and 3 indicate that sgoldstino production and decay may lead to observable signals at  $e^+e^-$  colliders, in particular at LEP. In the case of  $S$ , the most interesting signals correspond to  $e^+e^- \rightarrow \gamma S$ ,  $ZS$  or  $e^+e^-S$ , followed by the decay  $S \rightarrow gg$ . Similar considerations apply to  $P$ . All these signals would deserve a dedicated experimental analysis, including the comparison with the SM backgrounds. In the absence of a positive evidence, these analyses could be converted into stringent combined bounds on the gravitino and sgoldstino masses. While waiting for such an experimental study, we can only give a tentative picture of the LEP discovery potential for sgoldstinos, summarized in fig. 8. In drawing this picture, we just assumed some representative values for the relevant branching ratios and cross-sections. The choice of  $BR(Z \rightarrow S\gamma) = 10^{-5}$  can be partially justified on the basis of some existing OPAL and L3 studies [18], applicable in the mass region  $20 \text{ GeV} < m_S < 80 \text{ GeV}$ . As for the other processes, we are not aware of experimental studies whose results could be directly applied. In the case of  $\gamma S$  production, one should generalize the LEPI analyses of [18] along the lines of [19]. In the case of  $ZS$  production, one could exploit some similarities with SM Higgs searches (see, e.g., [20]

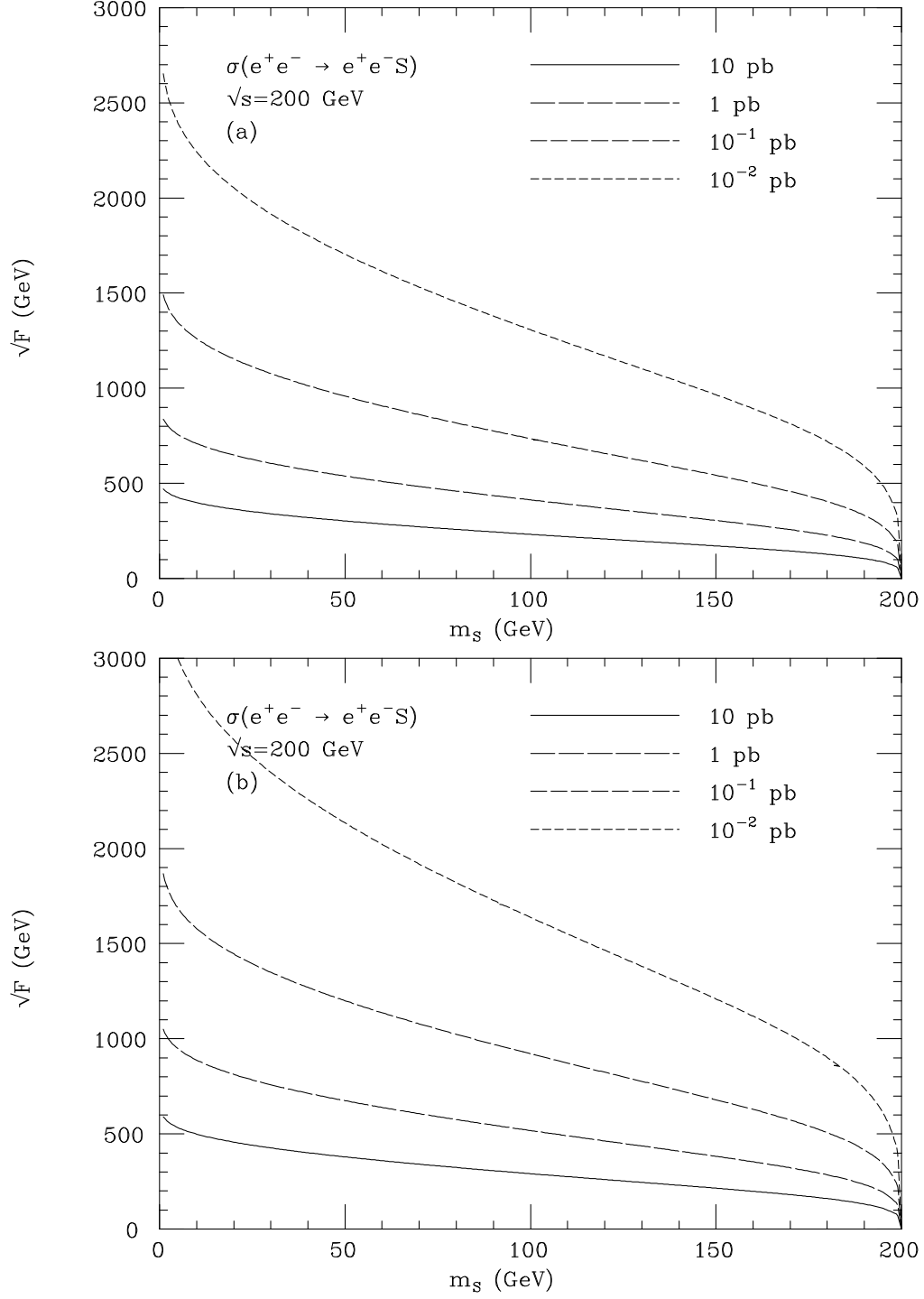


Figure 7: Lines corresponding to fixed values of  $\sigma(e^+e^- \rightarrow e^+e^-S)$  for  $\sqrt{s} = 200$  GeV, in the  $(m_S, \sqrt{F})$  plane and for the parameter choices of table 1.

and references therein), keeping in mind two very important differences: *i*) the production cross-section depends not only on  $m_S$ , but also on the mass parameters of table 1 and on  $\sqrt{F}$ , so the one-to-one correspondence between mass and production cross-section, valid for the SM Higgs, is in general lost; *ii*) massive sgoldstinos decay into gluon jets, not into  $b$ -jets, thus one cannot exploit all the machinery of  $b$ -tagging techniques. Finally, in the case of  $e^+e^- \rightarrow e^+e^-S$ , we expect both leptons to go undetected along the beam pipe in most cases, in accordance with the validity of the approximation used in the computation. Therefore, this signal will presumably suffer from a larger SM background than the previous ones, where the sgoldstino is produced in association with a photon or a  $Z$ . Also, we expect the LEP sensitivity to vary strongly with the sgoldstino masses, the most difficult region being the one with  $m_S \sim m_Z$ . Because of these problems, we limited ourselves to plotting contours of  $\sigma = 10^{-1}$  pb for the processes with a photon or a  $Z$  in the final state, and of  $\sigma = 1$  pb for the signal corresponding to  $\gamma\gamma$ -fusion, leaving a detailed analysis to our experimental colleagues.

We can see from fig. 8 that sgoldstino searches at LEP are likely to explore virgin land in the parameter space of models with a superlight gravitino. For example, the present limit on  $\sqrt{F}$ , for heavy sgoldstinos and MSSM sparticles, is only slightly above 200 GeV [3]. The associated production of MSSM sparticles and gravitinos is only relevant for masses of the MSSM sparticles smaller than the values assumed here. Indirect constraints from the muon anomalous magnetic moment [16], electroweak precision data [14] and anomalous four-fermion interactions [12] just give complementary and comparable bounds. Unitarity bounds just require the supersymmetry-breaking masses to be smaller than  $\mathcal{O}(2-3) \times \sqrt{F}$ , a condition which is comfortably fulfilled along most of our sensitivity contours: the only problematic regions are those very close ( $\lesssim 5$  GeV) to the boundary of the phase space for the process under consideration, which should therefore be excluded from the analyses and left for investigations at higher energies. Another fact emerging from fig. 8 is the complementarity of the different signals: their relative importance will depend not only on the experimental sensitivity, but also on the specific values of the gaugino and higgsino masses.

Finally, even if the present work is focused on sgoldstino signals at  $e^+e^-$  colliders, we would like to add a few comments on the possibility of producing sgoldstinos at hadron colliders. At leading order in the strong interactions, the dominant production mechanism for massive sgoldstinos should be gluon-gluon fusion, since direct couplings to quark-antiquark pairs are suppressed by the corresponding quark masses. For sufficiently heavy sgoldstinos, the resulting signal would be a peak in the di-jet invariant mass distribution. Such a signal is not present in the limit of negligible sgoldstino masses, and was therefore neglected in previous studies [7, 9]. Another possibility would be to look for sgoldstino production in association with an electroweak gauge boson ( $\gamma, Z, W$ ) or a jet. In the first case, the relevant partonic processes are  $q\bar{q} \rightarrow \gamma S$ ,  $q\bar{q} \rightarrow ZS$ ,  $q\bar{q}' \rightarrow WS$ , whose cross-sections are the obvious generalizations of those given here for  $e^+e^- \rightarrow \gamma S$  and  $e^+e^- \rightarrow ZS$ . In the second case, there are several diagrams that may contribute at

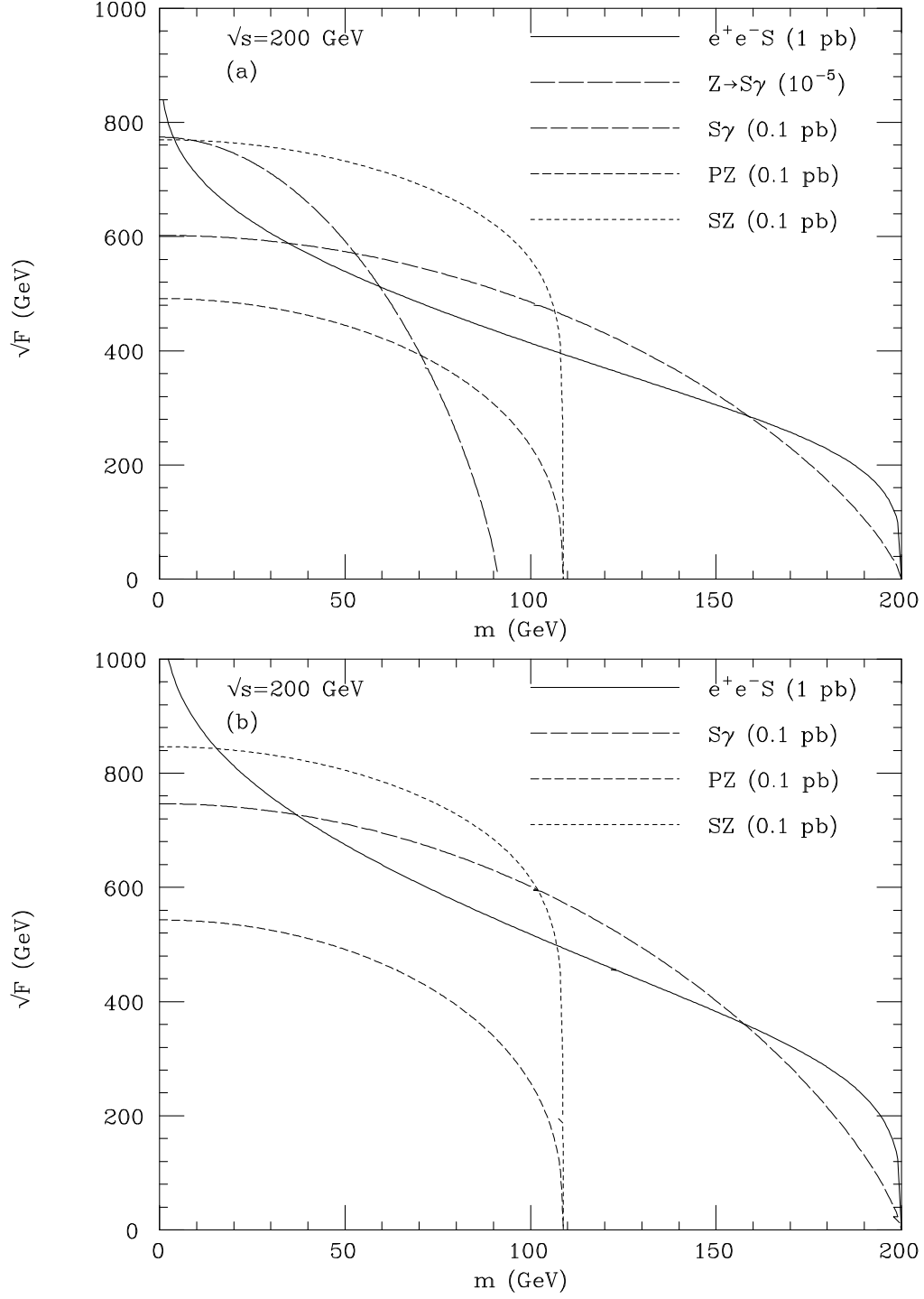


Figure 8: A tentative pictorial summary of the LEP discovery potential in the  $(m_{S(P)}, \sqrt{F})$  plane, for the parameter choices of table 1.

the same order in the strong interactions, thus the calculation of the cross-section is considerably more complicated than the existing ones [9], performed in the special case of negligible sgoldstino masses. We expect these processes to give constraints comparable with, and complementary to, the processes considered in the present paper. However, the theoretical complications due to the hadronic initial state and the presence of large SM backgrounds require a careful analysis: we are planning to come back to all this in a future publication.

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**Note added.** After the submission of the present paper, we were informed by P. Checchia and G. Wilson that, especially when the sgoldstino mass is close to  $m_Z$ , the study of the two-photon decay channel may lead to a sensitivity comparable with the two-gluon decay channel [21, 22].

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